

Left Right Model from Gauge Higgs Unification with Dark Matter

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We propose a five dimensional model based on the idea of Gauge Higgs Unification with the gauge group $SO(5) \times U(1)$ in Randall-Sundrum spacetime. We obtain a Left-Right symmetric model with a stable scalar therefore, a dark matter candidate. This scalar obtains a vacuum expectation value which gives mass to the top quark in the bulk through the Hosotani Mechanism. All the other fermions from the Standard Model are localized on a brane and obtain their mass from the vacuum expectation value of a brane localized scalar. This scalar fits the observed Higgs boson data. We are able to fit all the Standard Model observables while evading constraints.

I. INTRODUCTION

In the Standard Model (SM), the largest sources of free parameters are the set of Higgs to fermions couplings and the Higgs potential. We would like to understand how these parameters appear from a more fundamental theory.

One proposal for that is Gauge Higgs Unification (GHU), where one begins with an extra dimensional theory and the Higgs boson is the higher dimensional component of a gauge field [1–9]. In that case, the fermion-boson couplings as well as the scalar potential would be fixed by the gauge symmetry.

The most popular models begin with a gauge symmetry based on the group $SU(3)$ or $SO(5)$. These groups are able to generate a scalar with Higgs-like transformation properties under the SM group, but it can not accomodate its observables. To generate a realistic model one needs to make use of extra symmetries and fields [10]. One of such models is the so called Hosotani-Oda-Ohnuma-Sakamura (HOOS) model which is able to fit the measurements of the Higgs boson but requires many exotic fermions to do so [11, 12].

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The HOOS model has many interesting features which make it appealing. It has left-right (LR) symmetry that complies with electroweak precision tests. The scalar arising from the GHU idea, which we will call Hosotani scalar, is stable if one only has the SM fermionic content. The HOOS model adds many exotic fermions to fit the observed Higgs boson.

The Higgs boson isn't easily obtained from GHU, but we may accomodate another scalar field to come from the GHU idea. The scalar fields obtained this way may have special properties which may solve other problems, like dark matter [13].

We work with a different approach to HOOS model, in which instead of trying to get the SM directly, we obtain a LR model from it [14–17]. The Hosotani scalar is not the Higgs boson but one of the scalars that break the LR symmetry. This scalar is stable and a dark matter candidate. There are two more scalars which contribute to LR symmetry breaking, the Higgs boson comes from one of them. This boson has modified Higgs to boson couplings, with respect to the SM, that are within experimental constraints.

Except the top quark, all the fermions are localized on a brane and obtain their masses from the Higgs boson. The top quark propagates in the bulk and gets most of its mass from the Hosotani scalar.

II. $SO(5) \times U(1)$ MODEL IN THE BULK

The model is set up on the Randall - Sundrum (RS) warped spacetime with metric [18]:

$$ds^2 = e^{2\sigma(y)} \eta^{\mu\nu} dx_\mu dx_\nu + dy^2, \quad (1)$$

where $\sigma(y) = \sigma(y + 2L)$ and $\sigma(y) = k|y|$ for $|y| \leq L$.

The fundamental region for the fifth dimension is $0 \leq y \leq L$ and is limited by the UV brane at $y = 0$ and the IR brane at $y = L$.

The model is based on the gauge group $SO(5) \times U(1)$ and the corresponding gauge fields propagate in the bulk. The gauge group $SU(3)$ doesn't receive any special treatment.

After compactifying the fifth dimension we obtain the Kaluza-Klein (KK) mass scale [11]:

$$m_{KK} = \frac{\pi k}{e^{kL} - 1} \simeq \pi k e^{-kL} \quad (2)$$

that must comply with extra - dimensional constraints coming from electroweak precision tests [19].

In the bulk lives a fermion Ψ that transforms as $[\mathbf{5}, 2/3]$ under $SO(5) \times U(1)$. This fermion is needed to generate the scalar potential through the Hosotani mechanism [1, 20].

The action for this model is:

$$S = \int \sqrt{g} d^5x \left(i\bar{\Psi} \not{D} \Psi - c\sigma' \bar{\Psi} \Psi - \frac{1}{4} \text{Tr}[F^{MN} F_{MN}] - \frac{1}{4} B^{MN} B_{MN} \right), \quad (3)$$

where F^{MN} is the field strength of $SO(5)$, B^{MN} is the field strength of $U(1)$ and c is the kink mass of the fermion in the bulk.

A. Orbifolding, Hosotani phases and H parity

When we compactify, the components A_5 behave as scalars. We are free to choose the boundary conditions:

$$\begin{aligned} A_\mu(x, y_i - y) &= P_i A_\mu(x, y_i + y) P_i^{-1}, \\ A_5(x, y_i - y) &= -P_i A_5(x, y_i + y) P_i^{-1}, \\ \Psi(x, y_i - y) &= \eta^i \gamma^5 P_i \psi, \end{aligned} \quad (4)$$

where P_i are two 5×5 matrices, one for each brane, that obey $P_i^2, \eta^i = \pm 1$. The gauge symmetry on the branes will be reduced to the gauge transformations which satisfy:

$$P_i = \Omega(x, y_i - y) P_i \Omega^\dagger(x, y_i + y), \quad (5)$$

where Ω is a gauge transformation from $SO(5) \times U(1)$.

We choose $P_i = \text{diag}(1, 1, 1, 1, -1)$ and $\eta^i = 1$. These conditions break $SO(5) \times U(1) \rightarrow SO(4) \times U(1)$ on the branes. The four broken generators now behave as real scalars that transform as a $[\mathbf{4}, 0]$ under the remaining gauge group.

After breaking $SO(5)$ by orbifolding, the broken generators may obtain a vacuum expectation value (VEV) through the effective potential [21, 22]. We may align the VEV with a specific generator λ^H so the scalar that obtains a VEV can be written as:

$$A_5^H(x, y) = \theta_H(x) h(y) \lambda^H, \quad (6)$$

where $\theta_H(x)$ is a dimensionless scalar and $h(y)$ is its fifth dimensional profile, normalized as $\int_0^L h(y) dy = 1$.

Now let's study the transformation

$$\Omega(y) = e^{i\alpha \int_0^y h(y) \lambda^H} \quad (7)$$

that belongs to the broken generator of the group and changes $\theta_H \rightarrow \theta_H - \alpha$. We see that for specific values of α this may be a symmetry of the model if it satisfies eq.(5). In general this will be a symmetry for $\alpha = 2\pi$, hence the name of Hosotani Phase for θ_H . Since we are working only with tensorial representations of $SO(N)$, it is enough for $\alpha = \pi$ to be a symmetry of the system.

When we obtain the effective lagrangian for this scalar, namely [23]:

$$\mathcal{L}_{eff} = -V_{eff}(\theta_H) - m_W^2(\theta_H) W_\mu^\dagger W^\mu - \frac{1}{2} m_Z^2(\theta_H) Z_\mu Z^\mu - \sum_{a,b} m_{ab}^F(\theta_H) \bar{\psi}_a \psi_b. \quad (8)$$

The functions of θ_H have the symmetries:

$$\begin{aligned} V_{eff}(\theta_H + \pi) &= V_{eff}(\theta_H) = V_{eff}(-\theta_H), \\ m_{W,Z}^2(\theta_H + \pi) &= m_{W,Z}^2(\theta_H) = m_{W,Z}^2(-\theta_H), \\ m_{ab}^F(\theta_H + \pi) &= -m_{ab}^F(\theta_H) = m_{ab}^F(-\theta_H), \end{aligned} \quad (9)$$

due to the symmetry of the extra dimension and the fact that θ_H is a phase.

In this model with one heavy fermion in the bulk and the chosen gauge symmetry the VEV happens to be $\langle \theta_H \rangle = \pi/2$ [11, 24].

If we write $\theta_H = \frac{\pi}{2} + H_h/f$ where H_h is the scalar field and f is a dimensionful constant, the above mentioned functions of the effective lagrangian will behave like:

$$F\left(\frac{\pi}{2} + H_h/f\right) = F\left(\frac{\pi}{2} - H_h/f\right), \quad (10)$$

therefore they don't contain odd powers of H_h . This is called H-parity and makes the scalar field H_h stable [25, 26].

III. SCALARS IN THE MODEL

The IR brane has gauge symmetry $SO(4) \times U(1) \simeq SU(2)_R \times SU(2)_L \times U(1)_X$. Due to the RS metric, the dimensionful parameters are naturally small in this brane, solving the hierarchy problem. We will set every other field to live in this brane.

The broken generators behave as a scalar that transforms under the remaining gauge group as:

$$H_h \sim [\mathbf{2}, \mathbf{2}, 0] . \quad (11)$$

To break the LR symmetry completely we need two more scalars on the IR brane transforming like:

$$H_R \sim [\mathbf{2}, \mathbf{1}, 1] , \quad (12)$$

$$H_L \sim [\mathbf{1}, \mathbf{2}, 1] .$$

Each of these scalars has its usual symmetry breaking potential $V_{L,R} = \lambda_{L,R}(H_{L,R}^2 - v_{L,R}^2)^2$ that gives them VEVs v_L and v_R , breaking $SU(2)_{L,R} \times U(1)_X$ respectively.

There are stringent constraints on the masses for the W_R forcing $v_R \gtrsim 8$ TeV [19, 27, 28]. The left scalar obtains a VEV $v_L \sim O(100 GeV)$ so there are two orders of magnitude difference between these two VEVs. This should be tuned in the model.

Since we work in the low energy effective theory, where the KK modes are integrated out, the H_R scalar and the W_R^\pm, Z_R gauge bosons should be integrated out too, leaving us with the symmetry $SU(2)_L \times U(1)_Y$ on the IR brane with $Q_Y = T_R^3 + Q_X/2$. The scalars transform, under the remaining gauge group, as:

$$H_h \sim H_L \sim [\mathbf{2}, 1] \quad (13)$$

and effectively we have a Two Higgs Doublet Model (THDM).

The W boson receives its mass from the VEV of the Hosotani scalar in the bulk ($f\langle\theta_H\rangle$) and the left scalar VEV on the brane (v_L) so it looks like [2, 11, 24]:

$$S = \int \sqrt{g} d^5x w^+(y) w^-(y) \left(\frac{m_{KK}}{\pi L \sqrt{kL}} + \delta(y-L) \frac{g v_L}{2} \right) W^+(x) W^-(x), \quad (14)$$

where the $w(y)$ are the gauge boson's profiles in the fifth dimension. This means that the mass term that comes from the brane would have additional input from the profiles after normalizing the four-dimensional fields. Luckily the profiles of the W and Z in the fifth dimension are practically constant and there is not much input from them to the mass parameter nor to the charged currents [23] .

The VEV coming from the Hosotani mechanism ($f\langle\theta_H\rangle$) is aligned so that the mixing terms between W_R, Z_R and W_L, Z_L are practically zero [11, 17].

The masses for the W and Z are given just like in a THDM with $\rho = 1$ and the Weinberg angle has its usual value fixing the coupling constant of $U(1)_X$.

Due to H-parity there are no mixing terms between scalars. Since the H_h scalar breaks the symmetry in the bulk, it contains the Goldstone bosons so we work with it as:

$$H_h = \begin{pmatrix} 0 \\ \frac{v\frac{\pi}{2} + H(x)}{\sqrt{2}} \end{pmatrix}, \quad H_L = \begin{pmatrix} H^+ \\ \frac{v_L + h(x) + iA_0(x)}{\sqrt{2}} \end{pmatrix}, \quad (15)$$

where the extra $\pi/2$ comes from the Hosotani Mechanism and $H(x)$ is the remaining neutral scalar field in H_h (to be called Hosotani scalar too for the remaining of the paper).

Due to H-parity, $H(x)$ is stable and a viable candidate for dark matter [25].

Finally, the scalar potential is given by:

$$\begin{aligned} V = & -\mu_L^2 H_L^\dagger H_L + \frac{\lambda_L}{2} (H_L^\dagger H_L)^4 + \left\{ -\mu_{12}^2 H_L^\dagger H_h + h.c. \right. \\ & + \lambda_3 H_L^\dagger H_L H_h^\dagger H_h + \lambda_4 H_L^\dagger H_h H_h^\dagger H_L \\ & \left. + \frac{\lambda_5}{2} (H_L^\dagger H_h)^2 + \left(\frac{\lambda_6}{2} H_L^\dagger H_L + \frac{\lambda_7}{2} H_h^\dagger H_h \right) H_L^\dagger H_h + h.c. \right\} \Big|_H, \end{aligned} \quad (16)$$

where $\{\} |_H$ means that what is inside must respect H-parity, i.e. no odd power terms for $H(x)$.

IV. FERMIONS IN THE BULK AND BRANE

To have the full fermion content of the SM we need to add more fermions localized on the IR brane. If they were chosen to be in the bulk there would be many exotic fermions (since they would come from a bigger multiplet [11]). We will not talk about leptons since their behavior is the same as in a usual LR model.

First let's study the fermion in the bulk. Due to the chosen boundary conditions, the fermion multiplet breaks into:

$$\Psi \sim [\mathbf{5}, 2/3] \rightarrow P_L \oplus u_R^3 \sim [\mathbf{4}, 2/3] \oplus [\mathbf{1}, 2/3], \quad (17)$$

where the subindices indicate their chirality. Since $SO(4) \simeq SU(2)_L \times SU(2)_R$ the 4-tuplet fermion behaves as:

$$P_L \sim [\mathbf{2}, \mathbf{2}, 2/3] \quad (18)$$

under $SU(2)_R \times SU(2)_L \times U(1)_X$. Inside this bidoublet lies the third generation left quark doublet and an exotic doublet which will be heavy. The u_R^3 becomes the right-handed third generation up-type quark that gets mass through the Hosotani mechanism [11].

To complete the fermion content of the SM we add fermions in the brane with the following transformation properties under $SU(2)_R \times SU(2)_L \times U(1)_X$:

$$\begin{aligned} Q_L^a &= (u, d)_L^a \sim [\mathbf{1}, \mathbf{2}, 1/3], & Q_R^a &= (u, d)_R^a \sim [\mathbf{2}, \mathbf{1}, 1/3], \\ d_R^3 &\sim [\mathbf{1}, \mathbf{1}, -1/3], & \\ Q_R^4 &= (T, B)_R \sim [\mathbf{1}, \mathbf{2}, 5/3], \end{aligned} \tag{19}$$

where $a = 1, 2$, and the subindices L, R indicate their chirality. The leptons would behave as the light quarks.

The bidoublet P_L breaks into two doublets after the breaking of $SU(2)_R$. One of them has a TeV scale mass due to the term $y^4 Q_R^{4\dagger} P_L \tilde{H}_R$ and decouples from the theory. The other doublet behaves as a usual SM fermion doublet. We choose it to be the third family left fermion doublet since it is sufficiently heavy to generate the effective potential for the Hosotani scalar. The Hosotani scalar VEV gives mass to the top quark fixing the kink mass parameter c [11].

The Hosotani scalar obtains a VEV $\sim O(100 GeV)$ and it couples to the fermions on the IR brane giving them masses. The fact that the VEV comes from the bulk is important because it will be lowered in the IR brane when the IR fields absorb the warp factor (the same way it is done for the hierarchy problem) [18]. The fermion masses coming from the Hosotani scalar turn out to be $\sim O(10^{-13} GeV)$ and thus completely negligible.

In the low energy region, where $SU(2)_R$ is broken, we have an effective theory with cutoff $\Lambda \sim \text{TeV}$ and the terms $\frac{1}{\Lambda} Q_L H_L^\dagger Q_R H_R^\dagger$ become important. At these energies we have the same SM behavior of fermions and they all (except the top quark) get their masses due to the left scalar VEV v_L .

V. PHENOMENOLOGY

We choose $e^{kL} = 10^{15}$ so that the hierarchy problem is explained by this warp factor.

The main constraints for the extra dimensions parameters, in this model, come from KK mode's contributions to the S, T, U parameters [19, 20]. Since this model has a custodial symmetry and only one fermion in the bulk, the constraint is considerably lowered to [20]:

$$m_{KK} \gtrsim 850 \text{ GeV}.$$

This is a very low energy that would be in conflict with KK gluon constraints. Since this model doesn't contain $SU(3)$, that is not a problem, but if we were to include it, it should be confined to the IR brane.

This Hosotani scalar $H(x)$ obtains a mass which depends on m_{KK} [11, 24]. In figure 1 we see that dependence. We may also see that with the model we have there isn't a VEV for the Hosotani scalar if $m_{KK} > 1.7 \text{ TeV}$.

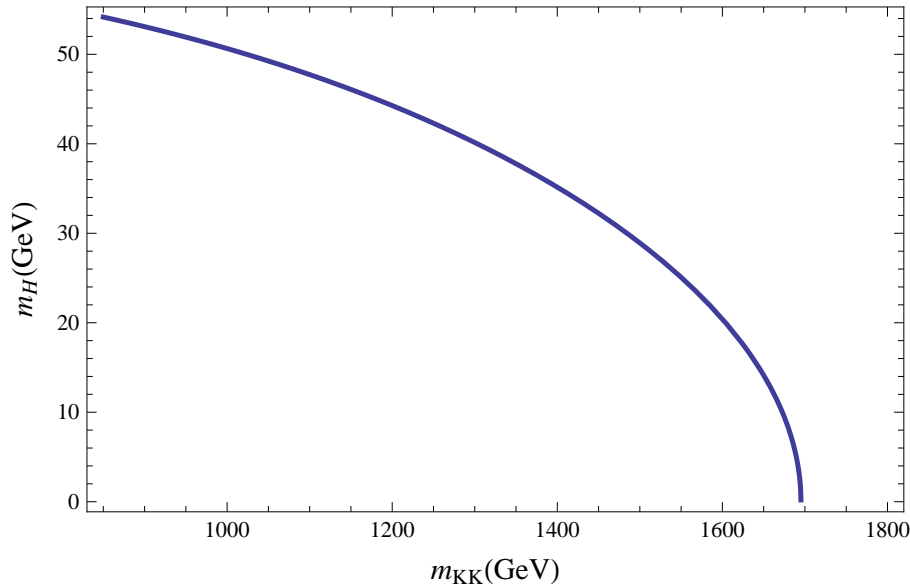


Figure 1: Mass of the Hosotani scalar.

As we see from the figure 1, the Hosotani scalar is light enough for the left scalar to decay into it with width:

$$\Gamma_{h \rightarrow HH} = \frac{1}{16\pi m_h} |6\lambda_6 \pi v + 2(2\lambda_3 + 2\lambda_4 + \lambda_5)v_L|^2 \sqrt{1 - 4m_H^2/m_h^2}. \quad (20)$$

The Hosotani scalar VEV depends on m_{KK} and since both scalars give masses to W and Z bosons then $\tan \beta = v/v_L$ depends on it too, as can be seen in figure 2.

The left scalar $h(x)$ should behave as the boson recently measured by ATLAS and CMS [29, 30]. The couplings of this scalar to fermions are SM like with extra an factor of $\cos \beta$,

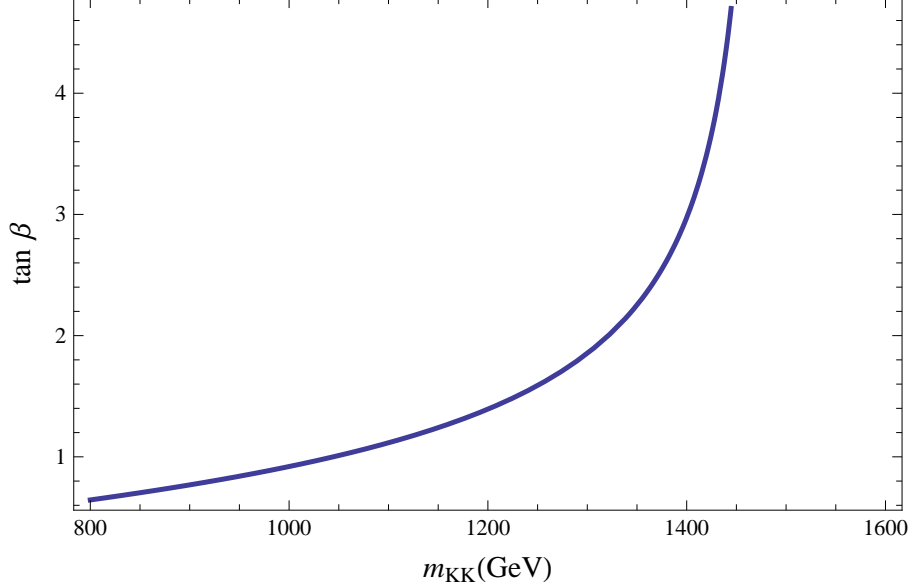


Figure 2: $\tan \beta$ dependence of extra dimensional parameters.

while its couplings to W and Z bosons have a factor of $1/\cos \beta$. As we see in figure 2, $\cos \beta$ decreases with m_{KK} . If we want the couplings of the left scalar to be SM like, then m_{KK} should be low. In our calculations we use $m_{KK} = 900$ GeV.

As we said above, the fermion masses coming from the Hosotani scalar are practically zero, except for the top. All other masses comes from v_L which resembles a type I THDM and therefore Flavor Changing Neutral Currents are negligible [31].

The scalar potential parameters define the masses of the scalars: λ_4 and λ_5 give the main contribution to the charged scalar and pseudoscalar masses, λ_6 for the left scalar mass and λ_7 for the Hosotani scalar. λ_L enters in all mass parameters but is not as important and the others. The remaining potential parameter λ_3 plays no important role in the masses but is significant in the $h - H$ coupling.

To obtain the branching ratios (BR) for the $h(x)$ we need to fix the potential parameters. We choose a representative set to give reasonable masses to the scalars: $v_L = 195$ GeV, $v = 150$ GeV, $\mu_{12} = 25$ GeV, $\lambda_4 = -0.676$, $\lambda_5 = -0.812$, $\lambda_6 = 0.050$, $\lambda_7 = 0.412$, $\lambda_L = 0.312$. The parameter λ_3 is left unfixed to explore the decay of the left scalar to the Hosotani scalar.

We obtain the masses for the scalar fields, with the parameters given above:

$$m_H = 53 \text{ GeV}, m_h = 125 \text{ GeV}, m_{A_0} = 342 \text{ GeV}, m_{H^\pm} = 271 \text{ GeV} \quad (21)$$

getting the mass for the Higgs-like particle and evading constraints for the other scalars [19].

In figure 3 we plot the $h(x)$ BRs varying the coupling of the scalars, that we call λ_h , and fixing the masses to the parameter set discussed above.

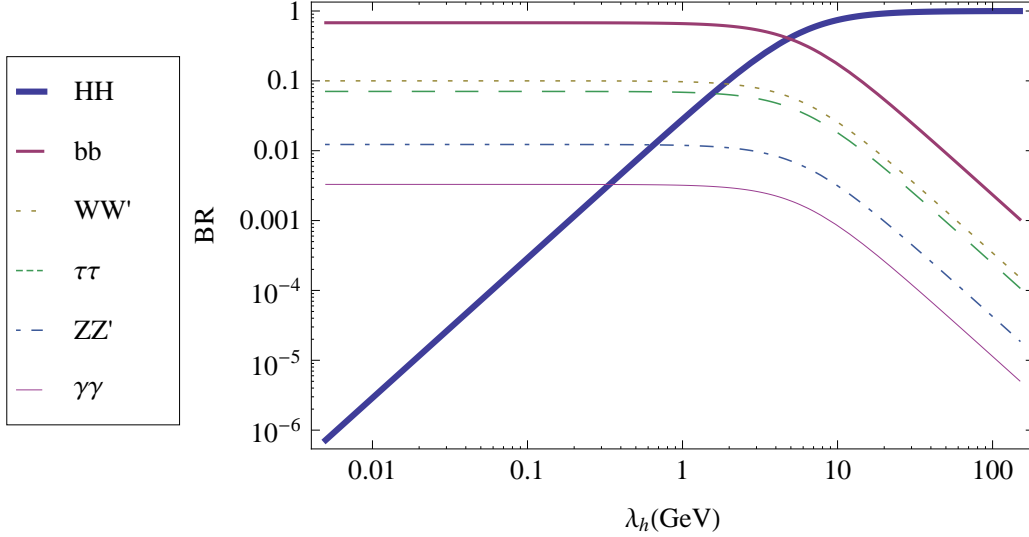


Figure 3: Branching ratios of the left scalar depending on the Hosotani-left scalar couplings.

We see in figure 3 that for a coupling $\lambda_h \sim O(100 GeV)$ (its biggest possible value assuming perturbativity) the decay into scalars is the main channel. This coupling can be tuned so that it almost vanishes and the channel is unnoticeable. In that case we obtain the BR for the W, Z, b compared to the SM ones as:

$$\sigma^{Z,W}/\sigma_{SM}^{W,Z} = 0.5, \quad \sigma^b/\sigma_{SM}^b = 1.2, \quad (22)$$

consistent with current measurements [29, 30].

The left scalar $h(x)$ gives mass to all the fermions (except the top) and fixes all its couplings but the top. The top quark receives its mass from the Hosotani scalar and the left scalar. The value of the top quark mass depends on its kink mass parameter c and its coupling to the left scalar. This leaves free the top - Higgs coupling y_t and this can take very small values (constrained by the measured Higgs production) up to the point of giving practically all of the top quark's mass (with this parameters $y_t = 0.88$). This is very interesting since very few models have a small Higgs - top coupling and would be important

in future experiments where this parameter can be measured such as future linear colliders and measurements of specific processes in the LHC [32, 33].

The decay of $h(x)$ to photons is enhanced by the charged scalar and the different couplings to fermions and W bosons, but it is affected by a change in the left scalar - top quark coupling. This coupling may be fixed to accommodate the diphoton channel measured value [29, 30]. With the chosen potential parameters and fixing the decay rate of the left scalar to photons to be equal as in the SM the Yukawa of the top quark is 0.65 the SM one.

VI. CONCLUSIONS

We have been able to build a realistic model that incorporates the idea of Gauge Higgs Unification but not in the usual sense. Gauge Higgs Unification proposes that the Higgs boson comes from the higher dimensional components of a gauge field. This idea needs too many extra fields to work, so in our model it is not the Higgs boson the one that comes from a GHU but another scalar.

We have a model with gauge symmetry based on the group $SO(5) \times U(1)$ and it lives in the RS spacetime. The gauge symmetry is broken to $SU(2)_L \times SU(2)_R \times U(1)_X$ via orbifolding, hence we have a LR symmetry that helps to comply with electroweak precision tests.

After compactification the higher dimensional components of the gauge fields behave as a scalar that obtains a VEV through the Hosotani mechanism. This Hosotani scalar is one of the scalar fields that break the LR symmetry but not the Higgs boson, as in usual GHU. Furthermore, due to symmetries of the model, it is stable and therefore a viable dark matter candidate.

There are two additional scalar field doublets under $SU(2)_{L,R}$ respectively that live on the IR brane and obtain a VEV, breaking the LR symmetry. The left doublet contains the Higgs boson and we fit its observables to the measured ones. The fact that these scalars live on the IR brane of the RS spacetime solve the hierarchy problem.

The $SU(2)_R \times U(1)_X$ symmetry breaks at the 10 TeV scale leaving effectively, at low energies, a THDM with $SU(2)_L \times U(1)_Y$ gauge symmetry, whose scalar sector contains a charged scalar and pseudoscalar within experimental constraints and the model is free of Flavor Changing Neutral Currents.

All the fermion content of the SM lives on the IR brane, except the top quark. The fermions on the brane obtain their masses from the left scalar VEV. The top quark comes from a multiplet in the bulk and receives its mass from a combination of both, the Hosotani mechanism and the left scalar VEV. Since its mass comes from two scalar VEVs the coupling of the left scalar field to the top quark is a free parameter. This result is unusual since we can have a small Higgs - top coupling and that would be interesting for future experiments.

We have been able to fit all the SM observables, including recent measurements from the Higgs boson. We evade all constraints from extra dimensions, left-right symmetries and extra fermions. The model is well constrained by this, but still viable.

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